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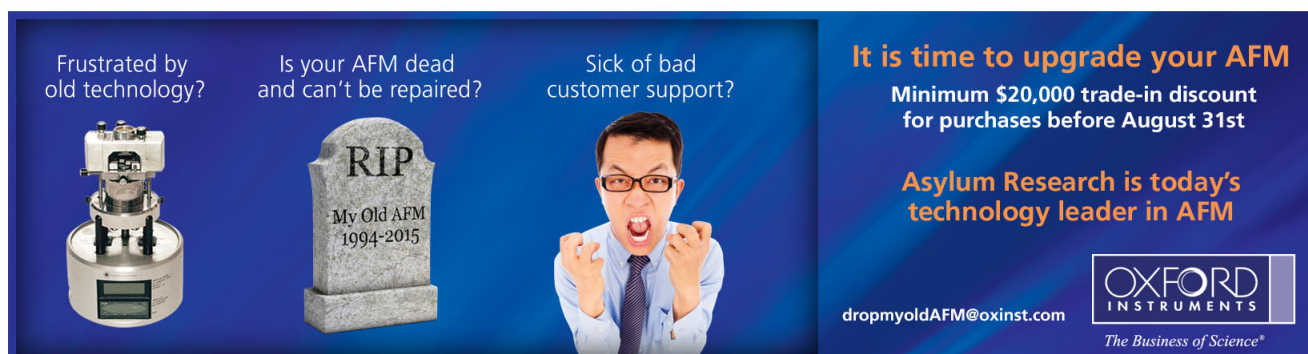
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Propagation of intense laser pulses in strongly magnetized plasmas

X. H. Yang,^{1,a)} W. Yu,^{2,b)} H. Xu,³ M. Y. Yu,^{4,5,c)} Z. Y. Ge,¹ B. B. Xu,¹ H. B. Zhuo,¹ Y. Y. Ma,¹ F. Q. Shao,¹ and M. Borghesi⁶

¹College of Science, National University of Defense Technology, Changsha 410073, China

²Shanghai Institute of Optics and Fine Mechanics, Chinese Academy of Sciences, Shanghai 201800, China

³School of Computer Science, National University of Defense Technology, Changsha 410073, China

⁴Institute for Fusion Theory and Simulation, Zhejiang University, Hangzhou 310027, China

⁵Institut für Theoretische Physik I, Ruhr-Universität Bochum, D-44780 Bochum, Germany

⁶School of Mathematics and Physics, Queen's University of Belfast, Belfast BT7 1NN, United Kingdom

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Propagation of intense circularly polarized laser pulses in strongly magnetized inhomogeneous plasmas is investigated. It is shown that a left-hand circularly polarized laser pulse propagating up the density gradient of the plasma along the magnetic field is reflected at the left-cutoff density. However, a right-hand circularly polarized laser can penetrate up the density gradient deep into the plasma without cutoff or resonance and turbulently heat the electrons trapped in its wake. Results from particle-in-cell simulations are in good agreement with that from the theory. © 2015 AIP Publishing LLC.

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The availability of intense short laser pulses has opened new areas of research on laser-plasma interactions, including basic physics as well as practical applications such as particle acceleration,^{1,2} fast ignition in inertial confinement fusion,^{3,4} and intense radiation sources at unusual wavelengths.^{5,6} Interaction of such lasers with plasma necessarily involves nonlinear effects, leading to processes such as self-focusing and filamentation,^{7–10} generation of solitons,^{7,11,12} electron and ion acceleration and heating,^{13,14} and plasma compression and shock generation.¹⁵ Intense short lasers can propagate into an unmagnetized inhomogeneous plasma only up to the relativistic cutoff density.

It is well known that the propagation of electromagnetic (EM) waves in plasma can be strongly affected by the presence of magnetic fields.^{16,17} Many earlier as well as current studies of EM waves in magnetized plasmas are on radio-wave propagation in the planetary ionospheres and atmospheres.^{16–21} Since several decades, the interaction between intense lasers and plasmas in the presence of weak magnetic fields has also been investigated.^{7,22,23} In the latter studies, the magnetic fields are usually limited to several to tens teslas and much lower than the wave magnetic field of the laser. On the other hand, extremely strong magnetic fields around or above several thousand teslas are ubiquitous in the cosmic environment²⁴ and can also occur in ultra-short ultra-intense laser interaction with plasmas because of high-current generation.^{25–28} In fact, an ultraintense magnetic field up to 75 000 T has been reported in laser interaction with dense plasmas.²⁹ It is therefore of interest to consider laser-plasma interaction in the presence of an external magnetic field with strength comparable to or larger than that of the laser light.

In this letter, we consider the parameter regime $a_0 \sim B_0$, where $a_0 = eE_0/mc\omega_0$, $B_0 = \omega_c/\omega_0$, $\omega_c = eB_{ext}/mc$ is the electron cyclotron frequency, ω_0 is the laser frequency, E_0 is

the laser electric field, $\mathbf{B}_0 = B_0\hat{z}$ is the external magnetic field in the laser propagation, or axial, direction z , c is the light speed, and e and m are the electron charge and rest mass, respectively. Accordingly, we consider the propagation of moderate ($a_0 \sim 0.74$) circularly polarized (CP) laser pulse into an inhomogeneous strongly magnetized ($B_0 = 2$, so that the laser frequency is far from the electron cyclotron frequency) plasma using particle-in-cell (PIC) simulation. It is found that, as expected, a left-hand CP laser is reflected at the left cutoff density $n_0 = \gamma_0 + B_0$, where n_0 is the electron density normalized by the critical density n_c . However, a right-hand CP laser can penetrate deeply into the plasma in the whistler wave mode without encountering cutoff or resonance.^{16,17} The plasma electrons in its turbulent wake can thus be continuously accelerated.

The motion of plasma electrons is governed by

$$d\mathbf{p} = \partial_t\mathbf{p} + (\mathbf{u} \cdot \nabla)\mathbf{p} = -\mathbf{E} - \mathbf{u} \times (\mathbf{B} + \mathbf{B}_0), \quad (1)$$

where $\gamma = \sqrt{1 + p^2}$ is the relativistic factor. The electric and magnetic fields \mathbf{E} and \mathbf{B} of the laser light, the electron velocity $u\hat{e}$, and the momentum $\mathbf{p} = \gamma u\hat{e}$ are normalized by $mc\omega_0/e$, c , and mc , respectively. From the Maxwell's equations, the propagation of EM waves is governed by

$$\nabla \times (\nabla \times \mathbf{E}) + \partial_t \mathbf{E} = \partial_t (n\mathbf{u}), \quad (2)$$

$$\nabla \cdot \mathbf{E} = Zn_i - n, \quad (3)$$

where n and n_i are the electron and ion densities normalized by n_c , and Z is the ion charge number. The ions are assumed to be fixed. The response of plasma electrons to the EM wave can be separated into fast (directly wave driven, marked with tilde) and slow (background or slow electron response to the relativistic ponderomotive force of the EM waves, indicated by subscript 0) components, namely, $n = n_0 + \tilde{n}$, $\mathbf{u} = \mathbf{u}_0 + \tilde{\mathbf{u}}$, $\mathbf{p} = \mathbf{p}_0 + \tilde{\mathbf{p}}$, $\gamma = \gamma_0 + \tilde{\gamma}$, $\mathbf{E} = \mathbf{E}_0 + \tilde{\mathbf{E}}$, and $\mathbf{B} = \mathbf{B}_0 + \tilde{\mathbf{B}}$. The fast components are then governed by

$$\gamma_0 \partial_t \tilde{\mathbf{u}} = -\tilde{\mathbf{E}} - \tilde{\mathbf{u}} \times \mathbf{B}_0, \quad (4)$$

^{a)}Electronic mail: xhyang@nudt.edu.cn

^{b)}Electronic mail: weiyu@mail.shnc.ac.cn

^{c)}Electronic mail: myyu@zju.edu.cn

$$\nabla \times (\nabla \times \tilde{\mathbf{E}}) + \partial_t \tilde{\mathbf{E}} = -(n_0/\gamma_0)(\tilde{\mathbf{E}} + \tilde{\mathbf{u}} \times \mathbf{B}_0), \quad (5)$$

$$\nabla \cdot \tilde{\mathbf{E}} = \tilde{n}. \quad (6)$$

Straightforward algebra leads to^{16,17}

$$\nabla^2 \mathbf{E}_t + \varepsilon_B \mathbf{E}_t = 0, \quad (7)$$

$$(1 - n_0/\gamma_0)E_z = 0, \quad (8)$$

where E_t and E_z are the transverse and longitudinal wave electric field components, $\varepsilon_B = 1 - n_0/(\gamma_0 - \sigma B_0)$ is the electron dielectric function, and $\gamma_0 = (1 - u_t^2 - u_z^2)^{-1/2}$. Note that $\tilde{\gamma} = 0$ for CP EM waves. For uniform plasmas, we have $\partial_{zz} E_t = -k_0^2 E_t$, where k_0 is the wave number of the EM wave normalized by c/ω_0 . The wave equation then leads to the dispersion relation $k_0^2 = \varepsilon_B$.

The propagation of CP lasers in magnetized plasmas strongly depends on ε_B . Most of the earlier works on laser-plasma interaction are for weak-field regimes, i.e., with $B_0 < 1$ or $\omega_c < \omega_0$. In this case, the propagation of CP lasers will be cut off when $\varepsilon_B = 0$, or at the cutoff density $n_0 = \gamma_0 + B_0$ for a left-hand CP laser ($\sigma = -1$), and $n_0 = \gamma_0 - B_0$ for a right-hand CP laser ($\sigma = 1$). We can see that in the strong-field regime (i.e., with $B_0 > 1$ or $\omega_c > \omega_0$), the right-hand CP laser can propagate in the plasma without cutoff, since the dielectric constant ε_B is greater than unity. In fact, it propagates as a relativistic whistler wave. Classical whistler waves have been well known for more than half a century in connection to radio wave propagation in weakly ionized gases such as the Earth's atmosphere, where the geomagnetic field is extremely weak. However, one can still have $\omega_c > \omega$ since the typical radio wave frequency ω is even smaller. Moreover, (8) describes external-field aligned electrostatic (ES) plasma oscillations that exist only at the cutoff density,^{16,20} but here the latter (namely, n_0/γ_0) is relativistically modified.

Numerical simulations are performed using the relativistic 2D3V PIC code LAPINE.^{30,31} The target, a hydrogen plasma, has an exponentially increasing density profile $n = 0.1 \exp[(z - 5\lambda_0)/12\lambda_0]$ and is initially located in $z = 5 - 65\lambda_0$,

where $\lambda_0 = 1\mu\text{m}$ is the laser wavelength, corresponding to a critical density of $n_c = 1.12 \times 10^{21} \text{ cm}^{-3}$. The initial temperature of the electrons and protons is 100 eV. The simulation box is $70\lambda_0 \times 30\lambda_0$ with 2800×1200 cells. Each cell contains 36 numerical macroparticles per species in the simulations. A (left- or right-hand) CP laser pulse with intensity $I_0 = 1.5 \times 10^{18} \text{ W/cm}^2$ is incident normally from the left boundary. Both the spatial and temporal profiles of the laser are Gaussian, with spot radius of $5\lambda_0$ and duration of $33.3T_0$ (FWHM), where $T_0 \sim 3.3 \text{ fs}$ is the laser cycle. Absorbing boundary conditions are used for the fields at the boundaries. The external magnetic field of strength $B_0 = 2$ is along the laser propagation direction.

We first consider a left-hand CP pulse. One can see in Fig. 1 for the axial Poynting vector S_z obtained from the simulations that the laser can propagate forward only up to $z \sim 46\lambda_0$, corresponding to the left cutoff density $n_0 = 1 + B_0 = 3$ (note that here $\gamma_0 \sim 1$), in good agreement with the theory. The laser pulse is reflected at this point, as can be seen in Fig. 1(c). Recall that without the external magnetic field, the laser pulse would be reflected (not shown) at $z \sim 33\lambda_0$, corresponding to the critical density $n = n_c$. Fig. 1(d) for S_z along the laser axis clearly exhibits the laser propagation, cut-off, and reflection. Accordingly, penetration of even a left-hand CP EM wave into a plasma can be enhanced by the presence of an external magnetic field since the left cutoff density is always larger than the critical density.¹⁷

Figure 2 is for the axial Poynting vector S_z of a right-hand CP laser pulse, with the other parameters the same as that for Fig. 1. As to be expected from the theory, we see that the right-hand CP laser pulse can propagate up the density gradient into the plasma without encountering cutoff, resonance, or reflection (which would correspond to $S_z < 0$). In fact, it is still propagating forward at $t = 120T_0$, reaching $z = 65\lambda_0$, where the electron density is $n \sim 15n_c$ (not shown). One can also see in Fig. 2 that in the wake there is a bubble-like region of weak or null S_z that contains (to be demonstrated below) energetic electrons. Fig. 2(d) shows that as the laser propagates deeper into the plasma, the entire

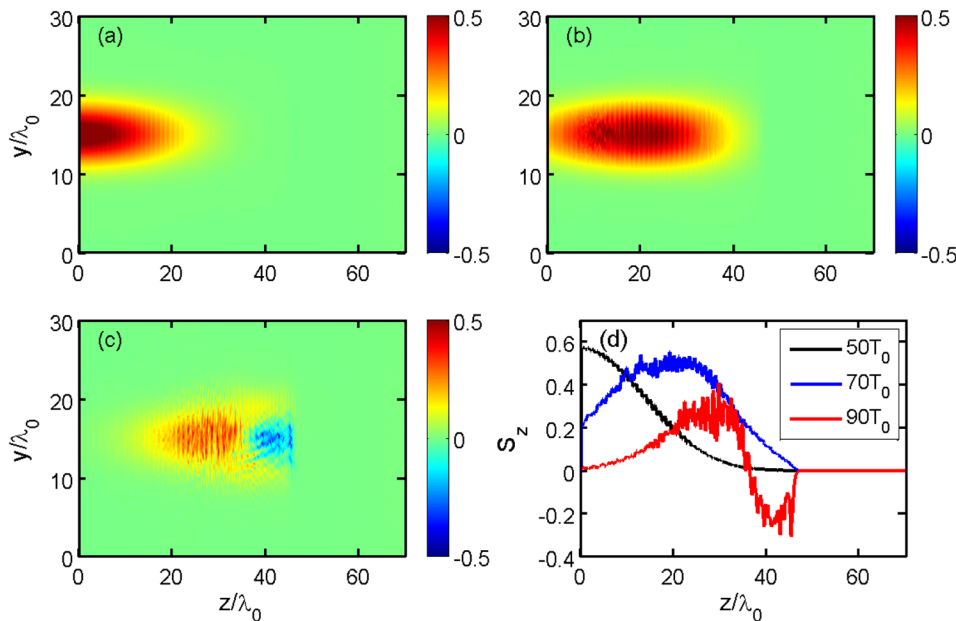


FIG. 1. The normalized (by $(mc\omega_0/e)^2$) axial Poynting vector S_z of a left-hand CP laser pulse at (a) $t = 50T_0$, (b) $70T_0$, and (c) $90T_0$, and (d) the Poynting vector S_z (averaged over a wavelength) at $y = 15\lambda_0$ along the axis. One can see that the laser propagation is cut off at $z/\lambda_0 \sim 46$.

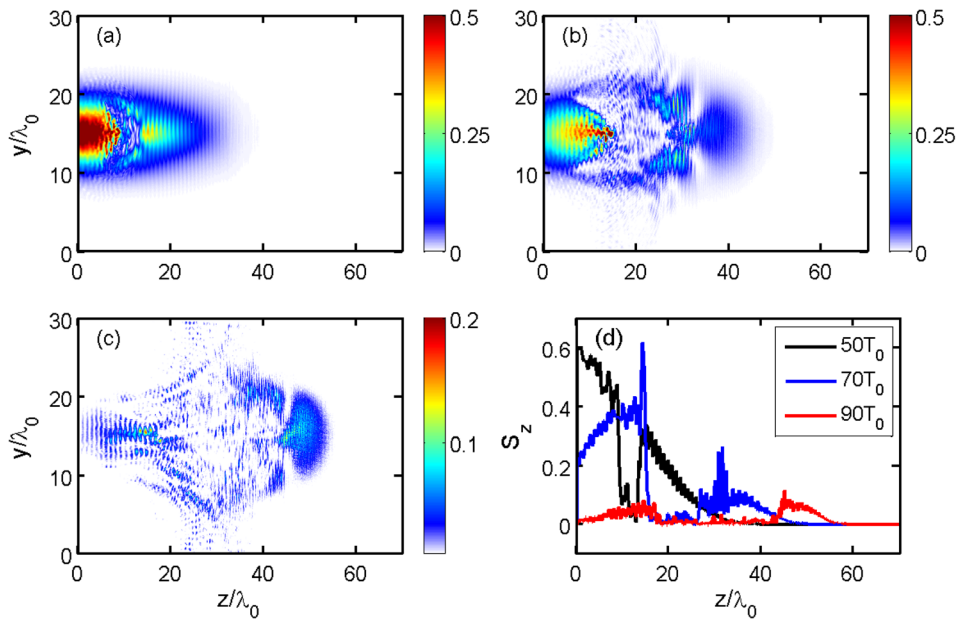


FIG. 2. The same as that in Fig. 1, but for a right-hand CP laser pulse.

wake field as well as this region expand. However, the laser pulse becomes weaker and weaker as it loses energy to the electrons continuously.

Figures 3(a) and 3(b) show the longitudinal electric field E_z and the axial momentum density $M_z = nP_z$. From their almost identical profiles, one can see that they correspond to the ES oscillations of the solution (8), which is limited to the region around $n_0/\gamma_0 \sim 1$, marked by a white rectangle in Fig. 3(d). Figs. 3(c) and 3(d) show the transverse and axial momentum of the electrons along the z direction. One can see that the electrons oscillate strongly in the wake region of the laser pulse. The maximum transverse and axial momentum of the electrons can reach $\sim 12mc$ and $\sim 25mc$, respectively, much higher than that of electrons accelerated by a laser in vacuum, where the electron momentum scales as $P_y = a_0 mc$ and $P_z = a_0^2 mc/2$.³² This result suggests that the plasma electrons can be much better accelerated by a right-hand CP laser.

The kinetic energy density K_e of electrons is shown in Figs. 4(a) and 4(b). In Figs. 4(a) and 4(b), we see that at $t = 70T_0$, the electrons in the wake region $z = 10 - 16\lambda_0$ are intensely heated, and at $t = 90T_0$ this region is much broadened (to $z = 27 - 34\lambda_0$) as the laser pulse propagates forward. Figs. 4(c) and 4(d) are for K_e and the transverse electromagnetic field energy W_t along the laser axis. We see that the profile of W_t is consistent with that of S_z in Fig. 2(d): as a result of energy transfer from the EM waves to the local electrons, the regions of high electron energy density almost coincide with that of low wave-energy density. Note that at $t = 90T_0$, the laser pulse has lost much of its energy, so that it becomes almost linear.

For comparison, in Fig. 5, we present the electron energy spectra for both left- and right-hand CP laser pulses, as well as that for a CP laser in unmagnetized plasma. Here, the average electron energy (or temperature) is 0.39 MeV for a CP laser propagation in unmagnetized plasma. A left-hand

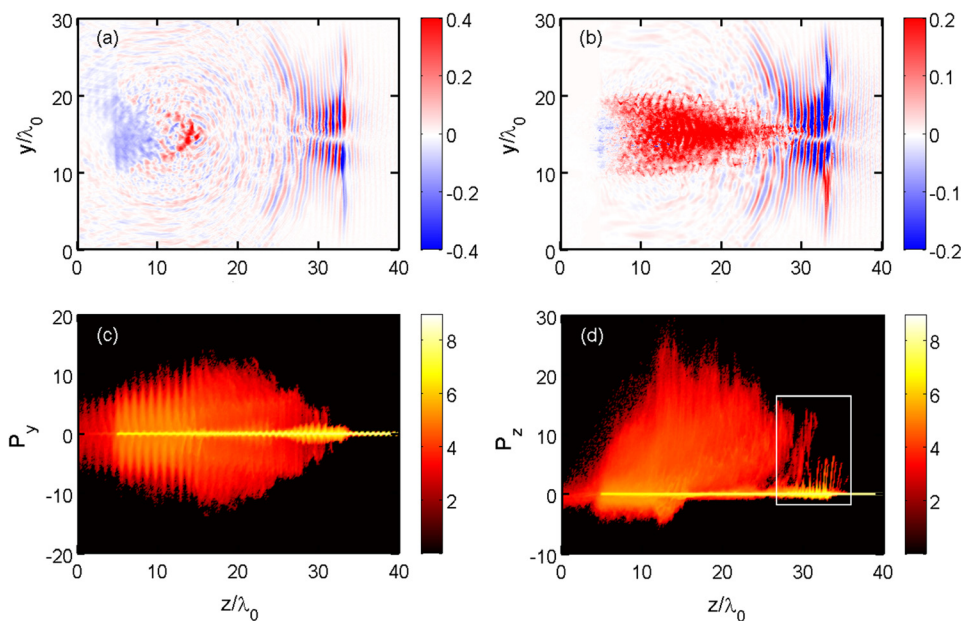


FIG. 3. Distribution of the longitudinal electric field E_z (a) and electron momentum density $M_z = nP_z$ (b) at $t = 70T_0$. (c) and (d) show the transverse and longitudinal electron momentum (on \log_{10} scale) P_y and P_z , respectively, at $t = 70T_0$. The momentum density and momentum are normalized by mcn_e and mc , respectively.

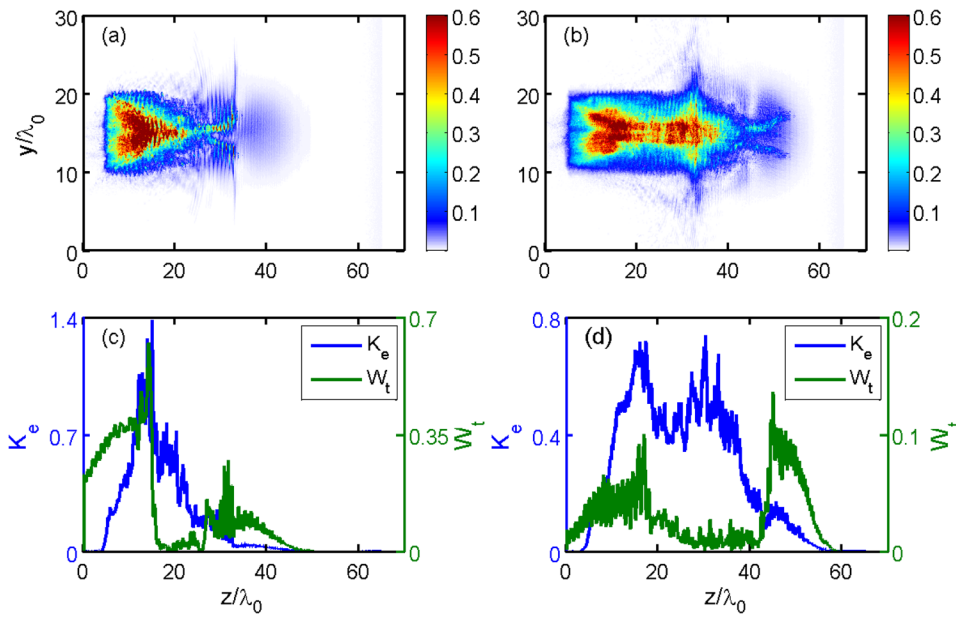


FIG. 4. The electron kinetic energy density K_e for a right-hand CP laser pulse at (a) $t = 70T_0$ and (b) $90T_0$. The corresponding profiles of (c) K_e and (d) transverse electromagnetic field energy W_t around $y = 15\lambda_0$ along the laser propagation direction, averaged over a wavelength. Both the electromagnetic-field and kinetic energy densities are normalized by $n_c mc^2$. Note the difference in the energy scales in (c) and (d).

CP laser can only increase it slightly to 0.56 MeV. Since its propagation is cut off at the left cutoff density, the acceleration as well as the number of electrons generated are limited and the energy cutoff is at about 4 MeV. On the other hand, with a right-hand CP laser pulse, the number of high energy electrons is significantly larger and the electron temperature can reach 2.7 MeV, with the cutoff energy at ~ 14 MeV. As mentioned, this occurs because a right-hand CP laser (with the dielectric constant ϵ_B larger than unity) can propagate in the plasma without cutoff and reflection. In the absence of an external magnetic field, a circularly polarized laser pulse is cut off at the critical density n_c , which is less than that of the left cutoff. However, since the EM wave is resonantly mode converted into the nonlinear ES electron plasma oscillations that are self trapped in the wake,⁷ the resulting electron acceleration/heating is more efficient than that of the left-hand CP waves (see Fig. 5).

In summary, we have shown that an intense right-hand CP laser pulse with frequency sufficiently less than the cyclotron frequency can propagate up the density gradient deep into an inhomogeneous strongly magnetized plasma as a

whistler wave without cutoff or resonance and continuously accelerate the plasma electrons trapped in its wake. This is possible because as the whistler wave of initial frequency $\omega_0 = \omega_c/2$ propagates up the density gradient, its frequency will be continuously downshifted,¹⁷ so that resonance becomes less and less likely even though relativistic and thermal effects can broaden the frequency domain for electron cyclotron resonance. In fact, one can see from the frequency spectrum $|E_\omega|$ (Fig. 6) of the transverse electric field (E_y) at $z = 30\lambda_0$ that besides the main peak corresponding to the laser pulse, there are only very weak peaks (from nonlinear harmonic generation), even though the second harmonic is close to ω_c . If there was cyclotron resonance, the energy spectrum would have a large peak near $\omega = \omega_c$. That is, the acceleration/heating of electrons found here is not directly by the whistler waves, which are often invoked for explaining phenomena in the geo and solar-corona magnetic fields.^{18–21} Instead, here the electron acceleration/heating takes place in the wake of the laser pulse, in a manner similar to that of classical wake acceleration in unmagnetized

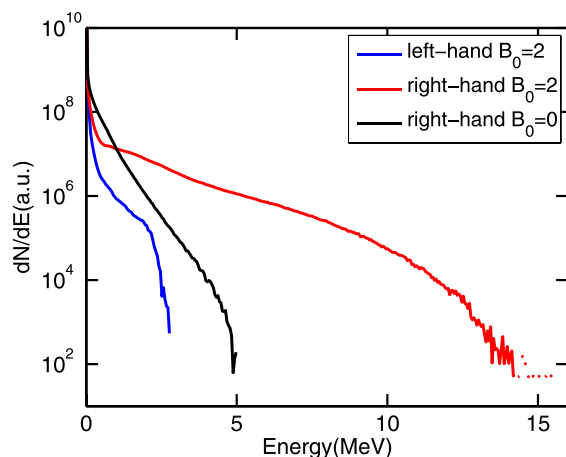


FIG. 5. Energy spectra of the electrons for left- and right-hand CP laser propagation in magnetized plasmas, as well as that in unmagnetized plasma.

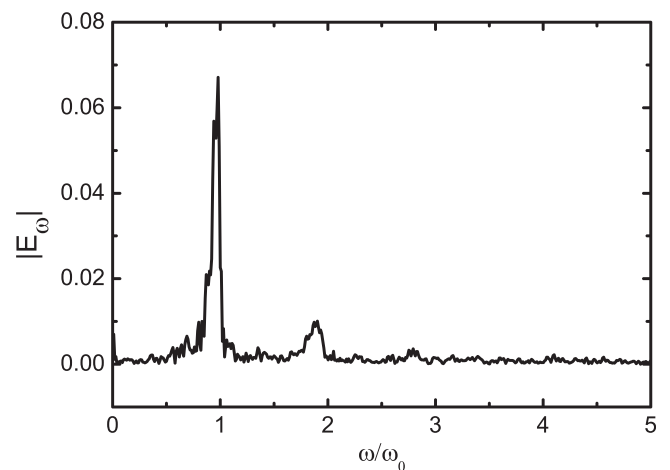


FIG. 6. Frequency spectrum $|E_\omega|$ (arbitrary units) of the transverse electric field (E_y) at $z = 30\lambda_0$. Note that $\omega_0 = \omega_c/2$, so that the main peak is at $\omega \leq \omega_c/2$.

plasmas.⁷ The present results should therefore be relevant to heating of geo and solar-corona plasmas without invoking electron cyclotron resonance, as well as to the fast ignition scheme in inertial confinement fusion,³³ target normal sheath acceleration of ions,^{34,35} as well as heating of magnetically confined plasmas.³⁶

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